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$\mathcal{B}(B_s \to \mu^+\mu^-)$ as an electroweak precision test

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ABSTRACT

Using an effective-theory approach, we analyze the impact of $\mathcal{B}(B_s \to \mu^+\mu^-)$ in constraining new-physics models that predict modifications of the Z-boson couplings to down-type quarks. Under motivated assumptions about the flavor structure of the effective theory, we show that the bounds presently derived from $\mathcal{B}(B_s \to \mu^+\mu^-)$ on the effective Z-boson couplings are comparable (in the case of minimal flavor violation) or significantly more stringent (in the case of generic partial compositeness) with respect to those derived from observables at the Z peak.

1. Introduction

The rare decay $B_s \to \mu^+\mu^-$ is one of the most clean low-energy probes of physics beyond the Standard Model (SM). A first experimental evidence of this rare process has recently been obtained by the LHCb Collaboration [1], that reported a 3.5σ signal. The corresponding flavor-averaged time-integrated branching ratio determined by LHCb is [1]

$$\overline{\mathcal{B}}^{\rm exp} = \left(3.2^{+1.5}_{-1.2}\right) \times 10^{-9}, \tag{1}$$

where the error is dominated by the statistical uncertainty and is expected to be improved significantly in the near future. At this level of precision there is good agreement with the SM prediction, that for the same quantity reads [2]

$$\overline{\mathcal{B}}^{\rm th}_{\rm SM} = (3.54 \pm 0.30) \times 10^{-9}, \tag{2}$$

taking into account the effect of $\Delta\Gamma_s \neq 0$ pointed out in Ref. [3].

The effectiveness of $B_s \to \mu^+\mu^-$ as a probe of physics beyond the SM is related to a double-suppression mechanism at work within the SM. One the one hand, it is a flavor-changing neutral-current (FCNC) process and, as such, it receives no tree-level contributions. On the other hand, the purely leptonic final state and the pseudoscalar nature of the initial state imply a strong helicity suppression and forbid photon-mediated amplitudes at the one-loop level. As a result of this double suppression, up to the one-loop level $B_s \to \mu^+\mu^-$ receives contributions only from Yukawa and weak interactions.

This process is often advocated as a probe of models with scalar-mediated FCNCs, that are naturally predicted in models with an extended Higgs sector. However, it is also an excellent probe of the $Z \to b\bar{s}$ effective coupling (see e.g. Refs. [4–6]). In this Letter we compare the bounds set on such coupling by $\mathcal{B}(B_s \to \mu^+\mu^-)$ with the deviations from universality on the $Z \to b\bar{b}$ coupling determined from electroweak precision observables. To this purpose, we describe the possible deviations on the Z-boson couplings to down-type quarks by means of an effective-theory approach, and we employ two motivated assumptions about the flavor structure of the theory, namely minimal flavor violation or generic partial compositeness, to relate flavor-changing and flavor-diagonal couplings.

2. Effective couplings of the Z boson to down-type quarks

As pointed out in Refs. [4,6], there exists a wide class of models where the only relevant deviations from the SM in $\mathcal{B}(B_s \to \mu^+\mu^-)$ and $Z \to b\bar{b}$ can be described in terms of modified Z-boson couplings at zero momentum transfer, defined by the following effective Lagrangian

$$\mathcal{L}^Z_{\rm eff} = \frac{g}{c_W} Z_\mu \bar{d}^i \gamma^\mu \left[\left(g^{ij}_L + \delta g^{ij}_L\right) P_L + \left(g^{ij}_R + \delta g^{ij}_R\right) P_R\right] d^j. \tag{3}$$

Here g is the $SU(2)_L$ gauge coupling, $c_W = \cos\theta_W$ ($s_W = \sin\theta_W$), and $g^{ij}_{L,R}$ denote the effective SM couplings. In the following we employ state-of-the-art expressions to estimate the SM contributions to $\mathcal{B}(B_s \to \mu^+\mu^-)$ and $Z \to b\bar{b}$, and use $\mathcal{L}^Z_{\rm eff}$ at the tree level only to estimate the non-standard effects parameterized by $\delta g^{ij}_{L,R}$.

For later convenience we recall the leading structure of the $g^{ij}_{L,R}$. The tree-level SM couplings are

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$$\left(g_L^{ii}\right)_{\text{tree}} = -\frac{1}{2} + \frac{1}{3}s_W^2, \qquad \left(g_R^{ii}\right)_{\text{tree}} = \frac{1}{3}s_W^2,$$
$$\left(g_{L,R}^{i\neq j}\right)_{\text{tree}} = 0. \tag{4}$$

At the one-loop level the $g_{L,R}^{ii}$ are gauge dependent, but they assume the following simple and gauge-independent form in the limit $m_t \gg m_W$ (or $g \to 0$):

$$\left(g_L^{ij}\right)_{\text{1-loop}}^{(g=0)} = \frac{m_t^2}{16\pi^2 v^2} V_{ti}^* V_{tj}, \qquad \left(g_R^{ij}\right)_{\text{1-loop}}^{(g=0)} = 0, \tag{5}$$

where V_{ij} denote the elements of the CKM matrix and $v \approx 246$ GeV.

The new-physics contributions, parameterized by $\delta g_{L,R}^{ij}$, can be related to the couplings of a manifestly gauge-invariant Lagrangian,

$$\mathcal{L}_{\text{eff}}^{\text{NP}} = -\frac{1}{2}\sum_{n,A}\sum_{i,j}\frac{c_{nA}^{ij}}{\Lambda^2}O_{nA}^{ij}, \tag{6}$$

with the following set of dimension-six operators:

$$\mathcal{O}_{1L}^{ij} = i\left(\overline{Q}_L^i \gamma^\mu Q_L^j\right)H^\dagger \overleftrightarrow{D}_\mu H, \qquad \mathcal{O}_{1R}^{ij} = i\left(\overline{D}_R^i \gamma^\mu D_R^j\right)H^\dagger \overleftrightarrow{D}_\mu H,$$
$$\mathcal{O}_{2L}^{ij} = i\left(\overline{Q}_L^i \tau^a \gamma^\mu Q_L^j\right)H^\dagger \tau^a \overleftrightarrow{D}_\mu H. \tag{7}$$

Defining the flavor indices $\{i,j\}$ in the mass-eigenstate basis of down-type quarks we find

$$\delta g_L^{ij} = \frac{v^2}{4\Lambda^2}\left(c_{1L}^{ij} + \frac{1}{4}c_{2L}^{ij}\right), \qquad \delta g_R^{ij} = \frac{v^2}{4\Lambda^2}c_{1R}^{ij}. \tag{8}$$

The set of operators in Eq. (7) is not the complete set of gauge-invariant dimension-six operators contributing to $B_s \to \mu^+\mu^-$ and $Z \to b\bar{b}$ at the tree level. In principle, we can consider also four-fermion (two-quarks/two-leptons) operators, terms of the type $J_\nu \times D_\mu F^{\mu\nu}$, or terms of the type $H^\dagger J_{\mu\nu} \times F^{\mu\nu}$, where J_ν and $J_{\mu\nu}$ are quark bilinears, and $F^{\mu\nu}$ generically denotes the field-strength tensor of $U(1)$ or $SU(2)_L$ gauge fields. However, the effects of these operators cannot be described by means of $\mathcal{L}_{\text{eff}}^Z$ and we lose the natural correlation between these two observables.[1] For this reason in the following we concentrate only on the set of operators in Eq. (7).

In order to relate flavor-diagonal and flavor-violating couplings we need to specify the flavor structure of the effective theory. We consider two reference frameworks: (1) the hypothesis of Minimal Flavor Violation (MFV), as defined in Ref. [7]; (2) the generic flavor structure implied by the hypothesis of Partial Compositeness (PC) [8], following the effective-theory approach described in Refs. [9,10].

In the MFV framework there is a strict correlation between flavor-diagonal (but non-universal) and flavor-violating couplings of the operators listed in Eq. (7). Restricting to the contributions relevant to this correlation, the effective couplings can be decomposed as follows:

$$\left(c_{nL}^{ij}\right)^{\text{MFV}} = a_{nL} \times \left(Y_u Y_u^\dagger\right)_{ij} \approx a_{nL}\frac{2m_t^2}{v^2}V_{ti}^* V_{tj}, \tag{9}$$

$$\left(c_{1R}^{ij}\right)^{\text{MFV}} = a_{1R} \times \left(Y_d^\dagger Y_u Y_u^\dagger Y_d\right)_{ij} \approx a_{1R}\frac{4m_{d_i}m_{d_j}m_t^2}{v^4}V_{ti}^* V_{tj}, \tag{10}$$

where $a_{nL,R}$ are unknown $O(1)$ couplings and $Y_{u,d}$ are the SM Yukawa couplings. The last equalities in Eqs. (9), (10) hold after rotating the Yukawa matrices in the mass-eigenstate basis of down-type quarks, where $Y_u = V^\dagger \lambda_u$ and $Y_d = \lambda_d$, with $\lambda_{u,d}$ diagonal matrices [7].

As a result, we can parameterize all the $\delta g_{L,R}^{ij}$ in terms of two flavor-blind parameters, $\delta g_{L,R}$, defined by

$$\left(\delta g_L^{ij}\right)^{\text{MFV}} = \frac{V_{ti}^* V_{tj}}{|V_{tb}|^2}\delta g_L,$$
$$\left(\delta g_R^{ij}\right)^{\text{MFV}} = \frac{m_{d_i}m_{d_j}}{m_b^2}\frac{V_{ti}^* V_{tj}}{|V_{tb}|^2}\delta g_R. \tag{11}$$

The normalization has been chosen such that

$$\delta g_{L(R)}^b \equiv \delta g_{L(R)}^{33} = \delta g_{L(R)}, \tag{12}$$

in order to identify $\delta g_{L,R}$ with the usual definition of the modified $Z \to b\bar{b}$ couplings [12]. As can be seen, in the left-handed sector the flavor structure is identical to the one of the leading one-loop contribution within the SM, reported in Eq. (5). In the right-handed sector the structure is different but the effects are expected to be very small due to the strong suppression of down-type masses. Indeed the overall normalization implies

$$\delta g_L^{\text{MFV}} = \frac{m_t^2|V_{tb}|^2}{2\Lambda^2}\left(a_{1L} + \frac{1}{4}a_{2L}\right), \qquad \delta g_R^{\text{MFV}} = \frac{m_b^2 m_t^2 |V_{tb}|^2}{v^2\Lambda^2}a_{1R}.$$

In the PC framework the correlation between flavor-diagonal and flavor-violating couplings is determined up to unknown $O(1)$ parameters, related to the hypothesis of flavor anarchy in the composite sector. In this case, following the notation of Ref. [10], we expect

$$\left(c_{nL}^{ij}\right)^{\text{PC}} \sim \frac{g_\rho^2\Lambda^2}{m_\rho^2}\epsilon_i^q\epsilon_j^q \propto |V_{ti}||V_{tj}|, \tag{13}$$

$$\left(c_{1R}^{ij}\right)^{\text{PC}} \sim \frac{g_\rho^2\Lambda^2}{m_\rho^2}\epsilon_i^d\epsilon_j^d \propto \frac{m_{d_i}m_{d_j}}{v^2|V_{ti}||V_{tj}|}, \tag{14}$$

where the $\epsilon_i^{q,d}$ parameterize the mixing of the SM fermions with the composite sector, and $\{m_\rho, g_\rho\}$ are the reference mass and coupling characterizing the composite sector. On the r.h.s. of Eqs. (13), (14) we have eliminated the $\epsilon_i^{q,d}$ in favor of quark masses and CKM angles by means of the relations [9,10]

$$\frac{|\epsilon_i^q|}{|\epsilon_j^q|} \sim \frac{|V_{ti}|}{|V_{tj}|}, \qquad \frac{|\epsilon_i^q\epsilon_i^d|}{|\epsilon_j^q\epsilon_j^d|} \sim \frac{m_{d_i}}{m_{d_j}}. \tag{15}$$

As can be seen, up to $O(1)$ factors the flavor structure of the left-handed couplings is the same as in the MFV framework. On the other hand, the structure is significantly different in the right-handed sector, where larger effects are now possible in the flavor-violating case. Ignoring $O(1)$ factors, we parameterize the structure of the two couplings in the PC framework as follows:

$$\left(\delta g_L^{ij}\right)^{\text{PC}} = \frac{|V_{ti}||V_{tj}|}{|V_{tb}|^2}\delta g_L,$$
$$\left(\delta g_R^{ij}\right)^{\text{PC}} = \frac{m_{d_i}m_{d_j}}{m_b^2}\frac{|V_{tb}|^2}{|V_{ti}||V_{tj}|}\delta g_R, \tag{16}$$

where again the normalization has been chosen in order to satisfy Eq. (12). (For recent studies of the same correlation within specific PC setups, see Ref. [11].) With such choice, the overall normalization implies

[1] The four-fermion operators do not contribute to $\mathcal{L}_{\text{eff}}^Z$ at the tree level, hence they have a negligible impact on $Z \to b\bar{b}$ compared to $B_s \to \mu^+\mu^-$. Conversely, operators with the field-strength tensor generate amplitudes suppressed by at least one power of p/v, with p the external momentum, that therefore have negligible impact on $B_s \to \mu^+\mu^-$ compared to $Z \to b\bar{b}$.

$$\delta g_L^{\rm PC} \sim \left(\frac{g_\rho \epsilon_3^q v}{2m_\rho}\right)^2, \qquad \delta g_R^{\rm PC} \sim \frac{1}{2}\left(\frac{m_b}{\epsilon_3^q m_\rho}\right)^2. \tag{17}$$

3. Analysis and discussion

The previous considerations can be summarized by stating that, within the two reference frameworks of MFV or PC, possible departures from the SM predictions in the $Z\bar{b}b$ couplings and in $\mathcal{B}(B_s \to \mu^+\mu^-)$ can be parameterized in terms of the two couplings $\delta g_{L,R}$ defined in Eq. (11) or Eq. (16).

Concerning Z-peak observables, the $\delta g_{L,R}$ shifts are constrained by R_b, A_b and $A_{\rm FB}^{0b}$. The state-of-the-art SM calculations for these quantities, to which it is straightforward to add the generic shifts in Eq. (12), can be implemented following Ref. [12] (taking also into account the recent SM estimate of R_b in Ref. [13]). These quantities can then be fitted to the averages of experimental results collected in Table 1, where we also report the main inputs necessary for their evaluation beyond the lowest order.

The resulting allowed regions at 68% CL and 95% CL in the δg_R–δg_L plane are shown in Fig. 1. As can be noticed, for both δg_L and δg_R the fit prefers positive non-zero values, and the SM point ($\delta g_R = \delta g_L = 0$) is outside the 95% CL region. The upper limits for the two parameters are

$$|\delta g_L|_{Z\bar{b}b} < 4.5 \times 10^{-3}, \qquad |\delta g_R|_{Z\bar{b}b} < 3.0 \times 10^{-2} \quad [95\%\ \text{CL}], \tag{18}$$

in good agreement with the results recently reported in Ref. [19].

Let us now compare these limits with those obtained from the $\mathcal{B}(B_s \to \mu^+\mu^-)$ measurement within the frameworks of MFV or PC. The $\delta g_{L,R}^{32}$ couplings shift linearly the Z-penguin contribution to the $\mathcal{B}(B_s \to \mu^+\mu^-)$ amplitude. These shifts can easily be translated into shifts on the short-distance function appearing in the SM formula for the branching ratio (see e.g. Ref. [2]). To good accuracy, the effect can simply be described by

$$\mathcal{B}\big(B_s \to \mu^+\mu^-\big) = \mathcal{B}\big(B_s \to \mu^+\mu^-\big)_{\rm SM} \times \left|1 + \frac{\sqrt{2}\pi^2}{G_F m_W^2 V_{tb}^* V_{ts}} \frac{(\delta g_L^{32} - \delta g_R^{32})}{Y_{\rm SM}}\right|^2, \tag{19}$$

where $Y_{\rm SM} \approx 0.957$.[2] Using the 95% CL range on the flavor-averaged branching ratio reported by LHCb [1]

$$1.1 \times 10^{-9} < \overline{\mathcal{B}}^{\rm exp} < 6.4 \times 10^{-9}, \tag{20}$$

and the central value of the SM prediction in Eq. (2) (at this level of accuracy the theoretical error is negligible), one obtains the following bounds on δg_L and δg_R:

$$\begin{aligned}
&|\delta g_L|^{\rm MFV,PC}_{B_s\to\mu^+\mu^-} < 2.3 \times 10^{-3},\\
&|\delta g_R|^{\rm MFV}_{B_s\to\mu^+\mu^-} < 1.0 \times 10^{-1},\\
&|\delta g_R|^{\rm PC}_{B_s\to\mu^+\mu^-} < 1.6 \times 10^{-4}.
\end{aligned} \tag{21}$$

These bounds have been obtained considering the effects of the two couplings separately (i.e. barring the possibility of cancellations between δg_L and δg_R, on which we will comment at the end of this section) and ignoring the fine-tuned configuration where the non-standard amplitude is about twice, and opposite in sign, compared to the SM one (a possibility that is highly disfavored by the $Z \to b\bar{b}$ constraints [6]).

Table 1
Input parameters relevant for the $Z \to b\bar{b}$ constraints. Quantities without an explicit reference are taken from Ref. [18]. We do not show the errors for quantities whose uncertainty has a negligible impact on our numerical analysis.

$M_h = 125$ GeV	[14]	$\Delta\alpha^{(5)}_{\rm had} = 0.02772$
$M_t = 173.2(0.9)$ GeV	[15]	$R_b = 0.21629(66)$
$\alpha_s(M_Z) = 0.1184(7)$	[16]	$A_b = 0.923(20)$
$\alpha^{-1}(M_Z) = 127.937$	[17]	$A^{0b}_{\rm FB} = 0.0992(16)$

These bounds are also depicted in Fig. 1 as horizontal or vertical bands delimited by solid lines. From the figure it is evident that, even with its large error, the recent evidence for $\mathcal{B}(B_s \to \mu^+\mu^-)$ provides a constraint on $|\delta g_L|$ – under either of the MFV or PC hypotheses – more stringent than the one obtained from the $Z \to b\bar{b}$ observables. Furthermore, the constraint on the $|\delta g_R|$ coupling within PC is stronger than the one obtained from the $Z \to b\bar{b}$ by more than two orders of magnitude. This circumstance is well represented by the right panel of Fig. 1, where the thickness of the $\mathcal{B}(B_s \to \mu^+\mu^-)$-allowed band (vertical blue 'line') is not resolved at the scale of the electroweak-fit ellipse. This implies that, within anarchic PC models, the $\mathcal{B}(B_s \to \mu^+\mu^-)$ bound forbids any significant contribution to $Z \to b\bar{b}$ observables able to decrease the existing tension between data and theoretical predictions.

As far as the bounds on the effective scale of new physics are concerned, in both frameworks the constraints derived from the $|\delta g_L|$ bound in Eq. (21) are largely dominant. They can be summarized as follows:

$$\begin{aligned}
&\Lambda > 2.6\ \text{TeV} \quad \big[\text{MFV}\,(\delta g_L)\big],\\
&m_\rho > \big(g_\rho \epsilon_3^q\big) \times 2.6\ \text{TeV} \quad \big[\text{PC}\,(\delta g_L)\big],
\end{aligned} \tag{22}$$

the equality of the numerical coefficient in the two cases being an accident due to the approximate relation $m_t|V_{tb}| \approx v/\sqrt{2}$. It is also worth mentioning the m_ρ bound implied by $|\delta g_R|$ in PC,

$$m_\rho > \frac{0.23\ \text{TeV}}{\epsilon_3^q} \quad \big[\text{PC}\,(\delta g_R)\big], \tag{23}$$

that becomes relevant in the limit $\epsilon_3^q \ll 1$, in which the bound from $|\delta g_L|$ gets weaker.

While the bounds in Eq. (21) are *per se* interesting, the present experimental error on $\mathcal{B}(B_s \to \mu^+\mu^-)$ does not do full justice to the sensitivity of this observable to possible modified Z-boson couplings. Therefore, we also considered the case of a $\mathcal{B}(B_s \to \mu^+\mu^-)$ measurement with central value as in Eq. (2) and error of $\pm 0.3 \times 10^{-9}$, that can be considered a realistic estimate of the experimental sensitivity on this observable around 2018. This statement takes into account the LHCb projections from Ref. [20], and the fact that CMS will likely produce a measurement with similar accuracy. We also assume a still subleading theoretical error, as expected by the steady progress in the lattice determination of the B_s decay constant [21]. With these assumptions on the projected total error on $\mathcal{B}(B_s \to \mu^+\mu^-)$, the 95% CL bounds on $\delta g_{L,R}$ become

$$\begin{aligned}
&|\delta g_L|^{\rm MFV,PC}_{[\sigma(B_s\to\mu\mu)=3\times10^{-10}]} < 4.6 \times 10^{-4},\\
&|\delta g_R|^{\rm MFV}_{[\sigma(B_s\to\mu\mu)=3\times10^{-10}]} < 2.0 \times 10^{-2},\\
&|\delta g_R|^{\rm PC}_{[\sigma(B_s\to\mu\mu)=3\times10^{-10}]} < 3.3 \times 10^{-5},
\end{aligned} \tag{24}$$

and the bounds in Eqs. (22) and (23) improve by a factor of about two. The comparison between Eq. (24) and Eq. (18) illustrates the

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[2] A similar expression holds for $\mathcal{B}(\bar{B}_s \to \mu^+\mu^-)$, with the replacement $(\delta g_L^{32} - \delta g_R^{32})/V_{tb}^* V_{ts} \to (\delta g_L^{23} - \delta g_R^{23})/V_{tb} V_{ts}^*$. Once $\delta g_{L,R}^{23,32}$ are expressed in terms of $\delta g_{L,R}$, the $\mathcal{B}(\bar{B}_s \to \mu^+\mu^-)$ and $\mathcal{B}(B_s \to \mu^+\mu^-)$ expressions are identical both in the MFV and in the PC parameterization, and can be directly compared with the flavor-averaged branching ratio reported by LHCb [1].

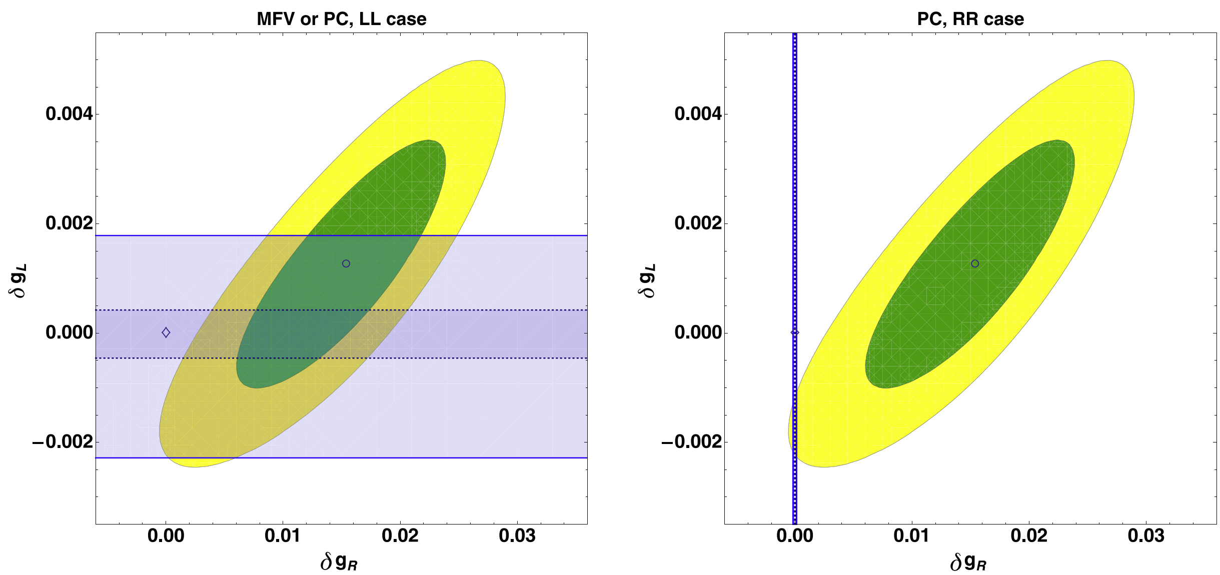

Fig. 1. Constraints on the couplings $\delta g_{L,R}$ describing the modified Z-boson couplings to down-type quarks. The inner and outer ellipses denote respectively the 68% and 95% CL regions as obtained from $Zb\bar{b}$ observables. The regions delimited by solid blue lines denote the 95% CL constraints from $\mathcal{B}(B_s \to \mu^+\mu^-)$ with the present precision, while those comprised between dotted lines are obtained with the $\mathcal{B}(B_s \to \mu^+\mu^-)$ accuracy expected by 2018 (see text for details). *Left panel:* δg_L constraint from $\mathcal{B}(B_s \to \mu^+\mu^-)$ under the hypotheses of either MFV or PC. *Right panel:* δg_R constraint from $\mathcal{B}(B_s \to \mu^+\mu^-)$ under the hypothesis of PC.

potential of uncovering even tiny new-physics deviations in the Z-boson couplings to down-type quarks via $\mathcal{B}(B_s \to \mu^+\mu^-)$. Note that, in the pessimistic case where no deviations from the SM prediction are observed in $\mathcal{B}(B_s \to \mu^+\mu^-)$, even the bound on δg_R within MFV will become more stringent compared to the one obtained from the $Z \to b\bar{b}$ observables.

Besides improvements in the $\mathcal{B}(B_s \to \mu^+\mu^-)$ measurement, a further avenue towards reducing the error on the $Z \to b\bar{s}$ effective coupling is, in principle, that of combining the constraints from other $b \to s$ decays, most notably $B \to K^*\mu^+\mu^-$ and $B \to K\mu^+\mu^-$ (recent attempts in this direction can be found in Ref. [22]; see also Ref. [23] for other related studies). However, the extraction of information about the $Z \to b\bar{s}$ effective coupling from these decays is not as pristine as in the $\mathcal{B}(B_s \to \mu^+\mu^-)$ case. In fact, on the one side, and at variance with $B_s \to \mu^+\mu^-$, these processes receive, already within the SM, substantial contributions from amplitudes other than the Z-penguin. In addition, the definition of observables related to these processes comes with inevitable theoretical assumptions, related to the dependence on additional hadronic form factors.

Finally, as anticipated, the bounds in Eq. (21) and Eq. (24) do not take into account the possibility of cancellations in the case where *both* δg_L and δg_R are switched on simultaneously. In practice, admitting such possibility does not lead to any significant changes in the plots of Fig. 1. As expected from the hierarchical nature of the bounds in Eqs. (21) or (24), the allowed region in the case of simultaneously non-zero δg_L and δg_R is dominated by the region allowed by the strongest constraint, namely δg_L in the case of MFV and δg_R in the case of PC.

4. Conclusions

The long-standing discrepancy between experimental data and SM predictions for the $Z \to b\bar{b}$ observables (A_{FB}^{0b} and, to a lesser extent, also R_b) has often been advocated as a possible hint of physics beyond the SM. If this is the case, under reasonable assumptions about the flavor structure of the new-physics model, sizable non-standard contributions should also be expected in $B_s \to \mu^+\mu^-$.

A first attempt to relate flavor-changing and flavor-diagonal constraints on the Z-boson couplings, under the assumption that they provide the dominant new-physics contribution to both $\mathcal{B}(B_s \to \mu^+\mu^-)$ and $Z \to b\bar{b}$, was made in Ref. [6]. At that time, the information from $Z \to b\bar{b}$ observables was used to derive possible upper bounds on $\mathcal{B}(B_s \to \mu^+\mu^-)$ and other FCNC processes. The situation is now reversed: the experimental precision reached on $\mathcal{B}(B_s \to \mu^+\mu^-)$ is such that this observable sets the dominant constraints on possible modified Z-boson couplings.

In MFV models, where sizable deviations are expected only in the left-handed couplings of the Z boson, the bound presently derived from $\mathcal{B}(B_s \to \mu^+\mu^-)$ is only slightly more stringent with respect to the one derived from $Z \to b\bar{b}$. However, the situation is likely to improve soon with the foreseen experimental progress on $\mathcal{B}(B_s \to \mu^+\mu^-)$, see Fig. 1 left. In generic models with partial compositeness, $\mathcal{B}(B_s \to \mu^+\mu^-)$ sets a constraint on possible modifications of the right-handed coupling considerably more stringent than $Z \to b\bar{b}$, see Fig. 1 right. This constraint forbids any significant contribution to $Z \to b\bar{b}$ observables able to decrease the existing tension between data and theoretical predictions.

More generally, our results illustrate how a measurement of $\mathcal{B}(B_s \to \mu^+\mu^-)$ with the expected accuracy of order 10% is able to unveil even tiny new-physics deviations in the Z-boson couplings to down-type quarks.

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References

[1] R. Aaij, et al., LHCb Collaboration, Phys. Rev. Lett. 110 (2013) 021801, arXiv:1211.2674.
[2] A.J. Buras, J. Girrbach, D. Guadagnoli, G. Isidori, Eur. Phys. J. C 72 (2012) 2172, arXiv:1208.0934.
[3] K. De Bruyn, R. Fleischer, R. Knegjens, P. Koppenburg, M. Merk, A. Pellegrino, N. Tuning, Phys. Rev. Lett. 109 (2012) 041801, arXiv:1204.1737 [hep-ph].

[4] M.S. Chanowitz, arXiv:hep-ph/9905478.
[5] G. Buchalla, G. Hiller, G. Isidori, Phys. Rev. D 63 (2000) 014015, arXiv:hep-ph/0006136.
[6] U. Haisch, A. Weiler, Phys. Rev. D 76 (2007) 074027, arXiv:0706.2054.
[7] G. D'Ambrosio, G.F. Giudice, G. Isidori, A. Strumia, Nucl. Phys. B 645 (2002) 155, arXiv:hep-ph/0207036.
[8] D.B. Kaplan, Nucl. Phys. B 365 (1991) 259;
T. Gherghetta, A. Pomarol, Nucl. Phys. B 586 (2000) 141, arXiv:hep-ph/0003129;
K. Agashe, G. Perez, A. Soni, Phys. Rev. D 71 (2005) 016002, arXiv:hep-ph/0408134;
R. Contino, T. Kramer, M. Son, R. Sundrum, JHEP 0705 (2007) 074, arXiv:hep-ph/0612180.
[9] S. Davidson, G. Isidori, S. Uhlig, Phys. Lett. B 663 (2008) 73, arXiv:0711.3376.
[10] B. Keren-Zur, P. Lodone, M. Nardecchia, D. Pappadopulo, R. Rattazzi, L. Vecchi, Nucl. Phys. B 867 (2013) 429, arXiv:1205.5803.
[11] R. Barbieri, D. Buttazzo, F. Sala, D.M. Straub, A. Tesi, arXiv:1211.5085 [hep-ph];
D.M. Straub, arXiv:1302.4651 [hep-ph].
[12] K. Hagiwara, Ann. Rev. Nucl. Part. Sci. 48 (1998) 463;
G.-C. Cho, K. Hagiwara, Y. Matsumoto, D. Nomura, JHEP 1111 (2011) 068, arXiv:1104.1769.
[13] A. Freitas, Y.-C. Huang, JHEP 1208 (2012) 050, arXiv:1205.0299 [hep-ph].
[14] ATLAS Collaboration, Phys. Lett. B 716 (2012) 1, arXiv:1207.7214;
CMS Collaboration, Phys. Lett. B 716 (2012) 30, arXiv:1207.7235.
[15] T. Aaltonen, et al., CDF and D0 Collaborations, Phys. Rev. D 86 (2012) 092003, arXiv:1207.1069.
[16] S. Bethke, Eur. Phys. J. C 64 (2009) 689, arXiv:0908.1135.
[17] LEP Electroweak Working Group, http://lepewwg.web.cern.ch.
[18] J. Beringer, et al., Particle Data Group Collaboration, Phys. Rev. D 86 (2012) 010001.
[19] B. Batell, S. Gori, L.-T. Wang, JHEP 1301 (2013) 139, arXiv:1209.6382.
[20] R. Aaij, et al., LHCb Collaboration, arXiv:1208.3355.
[21] See, e.g., E. Gamiz, arXiv:1212.4723;
C. Davies, PoS LATTICE 2011 (2011) 019, arXiv:1203.3862, and references therein.
[22] C. Bobeth, G. Hiller, D. van Dyk, C. Wacker, JHEP 1201 (2012) 107, arXiv:1111.2558;
F. Beaujean, C. Bobeth, D. van Dyk, C. Wacker, JHEP 1208 (2012) 030, arXiv:1205.1838;
D. Becirevic, N. Kosnik, F. Mescia, E. Schneider, Phys. Rev. D 86 (2012) 034034, arXiv:1205.5811;
W. Altmannshofer, D.M. Straub, JHEP 1208 (2012) 121, arXiv:1206.0273;
T. Hurth, F. Mahmoudi, Nucl. Phys. B 865 (2012) 461, arXiv:1207.0688;
S. Descotes-Genon, J. Matias, M. Ramon, J. Virto, JHEP 1301 (2013) 048, arXiv:1207.2753 [hep-ph].
[23] A. Behring, C. Gross, G. Hiller, S. Schacht, JHEP 1208 (2012) 152, arXiv:1205.1500 [hep-ph];
A.J. Buras, F. De Fazio, J. Girrbach, arXiv:1211.1896 [hep-ph].